

Drell-Yan Massive Lepton-Pair's Angular Distributions at Large Q_T

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Abstract

By measuring Drell-Yan massive lepton-pair's angular distributions, we can identify the polarization of the virtual photon of invariant mass Q which decays immediately into the lepton-pair. In terms of a modified QCD factorization formula for Drell-Yan process, which is valid even if $Q_T \gg Q$, we calculate the massive lepton-pair's angular distributions at large Q_T . We find that the virtual photons produced at high Q_T are more likely to be transversely polarized. We discuss the implications of this finding to the J/ψ mesons' polarization measured recently at Fermilab.

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Recent data on J/ψ polarization measured by CDF collaboration at Fermilab Tevatron seem to be inconsistent with the predictions from various models of J/ψ production [1]. The CDF data shows that the J/ψ mesons produced at $\sqrt{S} = 1.8$ TeV become more longitudinally polarized as the transverse momentum (Q_T) increases [1]. On the other hand, the various theoretical calculations predict the J/ψ mesons to be more transversely polarized at large Q_T [2].

In order to help solving this puzzle, it was proposed to measure the Drell-Yan massive lepton-pair's angular distributions at the kinematics similar to that of J/ψ production, and to extract the virtual photon's polarization from the lepton-pair's angular distributions [3,4]. Production of the virtual photons in the Drell-Yan process at large Q_T and small Q^2 has a lot in common with the production of J/ψ mesons at high Q_T . They both have two large physical scales: Q_T and Q^2 , which is equal to $M_{J/\psi}^2$ in the case of J/ψ production; and Q_T^2 is much larger than Q^2 . If the collision energy \sqrt{S} is large enough, and the logarithmic contributions from resumming all powers of $\ln^n(Q_T^2/Q^2)$ dominate the production cross sections, the virtual photon and J/ψ production will share the *same* partonic subprocesses, labeled by the R in Fig. 1. Then, the only difference between the virtual photon and J/ψ production at large \sqrt{S} and high Q_T is the difference in the fragmentation functions. The virtual photon fragmentation functions are completely perturbative if $Q^2 \gg \Lambda_{\text{QCD}}^2$ [4,5], while the parton to J/ψ fragmentation functions involve final-state nonperturbative soft interactions [6]. The measurements of the virtual photon and J/ψ polarization at high Q_T can help us to isolate the final-state effect in J/ψ production, and to narrow the questions about J/ψ formation.

In this letter, we calculate Drell-Yan massive lepton-pair's angular distributions in Quantum Chromodynamics (QCD). We present the polarization parameter α_{DY} as a function of Q_T , which can be in principle measured at Fermilab Tevatron. We find that the virtual photons at high Q_T are likely to be transversely polarized, which is different from the data on the polarization parameter $\alpha_{J/\psi}$ measured at Fermilab [1]. If our predictions for the α_{DY} are confirmed by future experiments, we may conclude that the non-perturbative final-state interactions are very important for the formation of J/ψ mesons.

When Q_T and Q^2 are both large, cross sections for Drell-Yan massive lepton-pair production in hadronic collisions, $A(P_A) + B(P_B) \rightarrow \gamma^*(\rightarrow l\bar{l}(Q)) + X$, can be systematically calculated according to QCD factorization theorem [7,8]. Since we are mainly interested in the cross sections at high Q_T and low Q^2 , we neglect the Z channel contributions in the following discussions.

After integrating over the azimuthal angle distribution, the Drell-Yan lepton-pair's angular distributions, as sketched in Fig. 2, can be expressed as [9]

$$\frac{d\sigma_{AB \rightarrow l\bar{l}(Q)+X}}{dQ^2 dQ_T^2 dy d\cos\theta} = \frac{d\sigma_{AB \rightarrow l\bar{l}(Q)+X}}{dQ^2 dQ_T^2 dy} \frac{3}{2(3 + \alpha_{\text{DY}})} [1 + \alpha_{\text{DY}} \cos^2\theta] \quad (1)$$

where the polarization parameter α_{DY} is given in terms of the cross sections for producing a polarized virtual photon,

$$\alpha_{\text{DY}} \equiv \left(\frac{d\sigma_{AB \rightarrow \gamma_T^*(Q)+X}}{dQ_T^2 dy} - \frac{d\sigma_{AB \rightarrow \gamma_L^*(Q)+X}}{dQ_T^2 dy} \right) / \left(\frac{d\sigma_{AB \rightarrow \gamma_T^*(Q)+X}}{dQ_T^2 dy} + \frac{d\sigma_{AB \rightarrow \gamma_L^*(Q)+X}}{dQ_T^2 dy} \right). \quad (2)$$

The subscripts, T and L in Eq. (2) represent the transverse and longitudinal polarizations, respectively. The inclusive Drell-Yan cross section in Eq. (1) can also be expressed in terms of the cross sections for producing a polarized virtual photon,

$$\frac{d\sigma_{AB \rightarrow \ell\bar{\ell}(Q)+X}}{dQ^2 dQ_T^2 dy} = \left(\frac{\alpha_{em}}{3\pi Q^2} \right) \left[2 \frac{d\sigma_{AB \rightarrow \gamma_T^*(Q)+X}}{dQ_T^2 dy} + \frac{d\sigma_{AB \rightarrow \gamma_L^*(Q)+X}}{dQ_T^2 dy} \right] \quad (3)$$

where the factor 2 is a consequence that the virtual photon has two transverse polarization states [9].

When Q_T and Q^2 are both large, the differential cross section for producing a virtual photon of polarization $\lambda = T$ or L in Eqs. (2) and (3) can be factorized as [7,8]

$$\frac{d\sigma_{AB \rightarrow \gamma_\lambda^*(Q)+X}}{dQ_T^2 dy} = \sum_{a,b} \int dx_1 \phi_{a/A}(x_1, \mu) \int dx_2 \phi_{b/B}(x_2, \mu) \frac{d\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*(Q)+X}}{dQ_T^2 dy}(x_1, x_2, Q, Q_T, y; \mu) \quad (4)$$

where $\sum_{a,b}$ run over all parton flavors, the μ represents both renormalization and factorization scales, and $\phi_{a/A}$ and $\phi_{b/B}$ are normal parton distributions. The short-distance parts, $d\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*(Q)+X}/dQ_T^2 dy$ in Eq. (4) are perturbatively calculable in terms of a power series in α_s [10]. When $Q_T \gg Q$, the perturbatively calculated short-distance parts receive a large logarithm $\ln(Q_T^2/Q^2)$ for every power of α_s beyond the leading order. Therefore, we need to resum such logarithms to all orders in α_s in order to derive a reliable cross section for the virtual photon production at large Q_T .

It was shown [11] that the large logarithms $\ln^n(Q_T^2/Q^2)$ come from the partonic subprocesses in which the virtual photon is produced from the decay of a parton that itself was produced at a distance scale $\sim 1/Q_T$. Because the parent parton can radiate massless partons before it radiates the virtual photon, there is one power of large $\ln(Q_T^2/Q^2)$ for an additional power of α_s . Since all powers of the large logarithms are from the final-state bremsstrahlung of a parton produced at a short-distance, such logarithms can be resummed into a fragmentation function for the parton to fragment into the virtual photon. Berger, Qiu, and Zhang [11] derived a modified factorization formula for calculating Drell-Yan cross section at $Q_T \geq Q$, which includes the resummation of the large logarithms,

$$\begin{aligned} \frac{d\sigma_{AB \rightarrow \gamma_\lambda^*(Q)+X}}{dQ_T^2 dy} &= \sum_{a,b,c} \int \frac{dz}{z^2} \int dx_1 \phi_{a/A}(x_1, \mu) \int dx_2 \phi_{b/B}(x_2, \mu) \\ &\quad \times \left[\frac{d\hat{\sigma}_{ab \rightarrow cX}}{dp_{cT}^2 dy} \left(x_1, x_2, p_c = \hat{Q}/z; \mu \right) \right] D_{c \rightarrow \gamma_\lambda^*}(z, \mu; Q) \\ &\quad + \sum_{a,b} \int dx_1 \phi_{a/A}(x_1, \mu) \int dx_2 \phi_{b/B}(x_2, \mu) \frac{d\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*(Q)+X}^Y}{dQ_T^2 dy}(x_1, x_2, Q, Q_T, y; \mu) \quad (5) \end{aligned}$$

where a, b , and c run over all parton flavors, \hat{Q}^μ is the on-shell part of momentum Q^μ , and μ represents the renormalization, factorization, and fragmentation scales. The first term on the right-hand-side of Eq. (5) represents the resummed contributions, which include all powers of the large logarithms; and the second term is free of the large logarithms and carries the same physical meaning as the Y -term in the resummation formula for $Q_T \ll Q$ [12]. In Eq. (5), the $D_{c \rightarrow \gamma_\lambda^*}$ are the fragmentation functions for the partons of flavor c to fragment into a virtual

photon of invariant mass Q and polarization λ , and their operator definitions are given in Ref. [4]. The resummation of the large $\ln(Q_T^2/Q^2)$ to all orders in α_s can be achieved by solving the evolution (or renormalization group) equations for these fragmentation functions [4]. In Eq. (5), the partonic part $d\hat{\sigma}_{ab \rightarrow cX}/dp_{cT}^2 dy$ are perturbatively calculable and represent the production of partons of flavor c at a distance scale $\sim 1/p_{cT} \sim 1/Q_T$ [13]. Since the Y -term is defined to be the perturbative difference between Drell-Yan cross section and the resummed part of the cross section, the corresponding partonic parts at the order of α_s^n are given by [11]

$$\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*}^{Y(n)}(Q)X \equiv \hat{\sigma}_{ab \rightarrow \gamma_\lambda^*}^{(n)}(Q)X - \sum_{m=2}^n \hat{\sigma}_{ab \rightarrow cX}^{(m)} \otimes D_{c \rightarrow \gamma_\lambda^*}^{(n-m)}(Q) \quad (6)$$

where \otimes represents the convolution over z as defined in Eq. (5), and $D^{(n)}$ are perturbatively calculated virtual photon fragmentation functions at order of α_s^n .

The modified factorization formula in Eq. (5) should be valid for all $Q_T \geq Q$. When $Q_T \gg Q$, the resummed part dominates the cross sections, and the Y -term represents a small correction. When $Q_T \sim Q$, the logarithms are small and the cross sections are dominated by the Y -term. All short-distance parts in Eq. (5) are evaluated at the same short-distance scale $\sim 1/Q_T$ and free of the large logarithms. All leading logarithmic contributions from a distance scale of $1/Q_T$ to $1/Q$ are resummed into the fragmentation functions. With the modified factorization formula in Eq. (5), we are able to calculate the polarization parameter α_{DY} in Eq. (2) reliably in QCD perturbation theory by calculating the short-distance parts order-by-order in α_s .

From the modified factorization formula in Eq. (5), we obtain the leading order contributions to the virtual photon cross sections by calculating the lowest order contributions to both $\hat{\sigma}_{ab \rightarrow c}$ and $\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*}^Y$. From Eq. (5), the lowest order contributions to the short-distance $d\hat{\sigma}_{ab \rightarrow cX}/dp_{cT}^2 dy$ are independent of the virtual photon's polarization, and are given by the lowest order two-to-two partonic tree diagrams at $O(\alpha_s^2)$,

$$\frac{d\hat{\sigma}_{ab \rightarrow cd}}{dp_{cT}^2 dy} (x_1, x_2, p_c = \hat{Q}/z; \mu) = \pi \frac{\alpha_s^2(\mu)}{x_1 x_2 S} \left| \frac{1}{g^2} \overline{M}_{ab \rightarrow cd} \right|^2 \delta(\hat{s} + \hat{t} + \hat{u}), \quad (7)$$

where g is the strong coupling constant and $S = (P_A + P_B)^2$ is the collision energy. In Eq. (7), \hat{s} , \hat{t} , and \hat{u} are the parton level Mandelstam variables. The lowest order two-to-two matrix element squares in Eq. (7) are given in Ref. [13].

The lowest order contributions to the partonic cross sections $\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*}$ in Eq. (6) are at order of $O(\alpha_{em}\alpha_s)$ from the subprocesses: $q\bar{q} \rightarrow \gamma^*(Q)g$ and $qg \rightarrow \gamma^*(Q)q$. Since the subtraction term in Eq. (6) starts at order of $O(\alpha_{em}\alpha_s^2)$, we have the leading order contributions to the Y -term as

$$\begin{aligned} \frac{d\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*}^{(Y-LO)}(Q)}{dQ_T^2 dy} &= \frac{d\hat{\sigma}_{ab \rightarrow \gamma_\lambda^*}^{(LO)}(Q)}{dQ_T^2 dy} \\ &= \pi \frac{\alpha_{em}(\mu)\alpha_s(\mu)}{x_1 x_2 S} \left| \frac{1}{eg} \overline{M}_{ab \rightarrow \gamma_\mu^* d} \epsilon_\lambda^\mu(Q) \right|^2 \delta(\hat{s} + \hat{t} + \hat{u} - Q^2), \end{aligned} \quad (8)$$

where e is the electron charge and \hat{s} , \hat{t} , and \hat{u} are again the parton level Mandelstam variables. The matrix element squares in Eq. (8) for producing a polarized virtual photon depend on the photon's polarization vector $\epsilon_\lambda^\mu(Q)$.

The choice of the polarization vectors are not unique [9]. But, it should be consistent with the definition of α_{DY} in Eqs. (1) and (2). The angle θ in Eq. (1) depends on the frame choice and its z -axis [9]. As a result, the extracted value of α_{DY} as well as the polarization vectors used to define the polarized virtual photon cross sections in Eq. (2) are also sensitive to the frame choice and its z -axis. For the consistency of the QCD factorization formula in Eq. (5), it is very important to use the same polarization vectors to calculate both the virtual photon fragmentation functions and the Y -term. Since we are interested in the region where $Q_T \gg Q$, we will not choose the Collins-Soper frame [14], which offers special advantage when $Q_T \ll Q$. Instead, as shown in Fig. 2, we choose the z -axis along the direction of the virtual photon's momentum \vec{Q} and angle θ to be the angle between the \vec{Q} and \vec{l} , the direction of one of the decay leptons. Our choice of the z -axis corresponds to the ‘‘S-helicity’’ frame defined in Ref. [9]. With our choice of the helicity frame and angle θ , $\alpha_{\text{DY}} = +1(-1)$ corresponds to the virtual photon in a purely transverse (longitudinal) polarization state.

To minimize the frame dependence, we present our results in terms of the covariant hadronic tensors $w_{[ab]\mu\nu}$ and polarization vector $\epsilon_\lambda^\mu(Q)$, which are defined as

$$\left| \frac{1}{eg} \overline{M}_{ab \rightarrow \gamma_\mu^* d} \epsilon_\lambda^\mu(Q) \right|^2 \equiv w_{[ab]\mu\nu} \epsilon_\lambda^{*\mu}(Q) \epsilon_\lambda^\nu(Q). \quad (9)$$

For quark-antiquark annihilation subprocess: $q(p_1)\bar{q}(p_2) \rightarrow \gamma^*(Q)g$, we obtain the hadronic tensor,

$$\begin{aligned} w_{[q\bar{q}] }^{\mu\nu} = e_q^2 \left(\frac{4}{9} \right) \left[\frac{1}{\hat{t}\hat{u}} \right] \left\{ -4Q^2 \left[\left(p_1^\mu - \frac{p_1 \cdot Q}{Q^2} Q^\mu \right) \left(p_1^\nu - \frac{p_1 \cdot Q}{Q^2} Q^\nu \right) \right. \right. \\ \left. \left. + \left(p_2^\mu - \frac{p_2 \cdot Q}{Q^2} Q^\mu \right) \left(p_2^\nu - \frac{p_2 \cdot Q}{Q^2} Q^\nu \right) \right] \right. \\ \left. - \left(g^{\mu\nu} - \frac{Q^\mu Q^\nu}{Q^2} \right) \left[(2p_1 \cdot Q)^2 + (2p_2 \cdot Q)^2 \right] \right\} \end{aligned} \quad (10)$$

with quark's fractional charge e_q ; and for the ‘‘Compton’’ subprocess: $g(p_1)q(p_2) \rightarrow \gamma^*(Q)q$, we have

$$\begin{aligned} w_{[gq] }^{\mu\nu} = e_q^2 \left(\frac{1}{6} \right) \left[\frac{1}{\hat{s}(-\hat{u})} \right] \left\{ -4Q^2 \left(p_1^\mu - \frac{p_1 \cdot Q}{Q^2} Q^\mu \right) \left(p_1^\nu - \frac{p_1 \cdot Q}{Q^2} Q^\nu \right) \right. \\ - 8Q^2 \left(p_2^\mu - \frac{p_2 \cdot Q}{Q^2} Q^\mu \right) \left(p_2^\nu - \frac{p_2 \cdot Q}{Q^2} Q^\nu \right) \\ - 4Q^2 \left[\left(p_1^\mu - \frac{p_1 \cdot Q}{Q^2} Q^\mu \right) \left(p_2^\nu - \frac{p_2 \cdot Q}{Q^2} Q^\nu \right) \right. \\ \left. \left. + \left(p_2^\mu - \frac{p_2 \cdot Q}{Q^2} Q^\mu \right) \left(p_1^\nu - \frac{p_1 \cdot Q}{Q^2} Q^\nu \right) \right] \right. \\ \left. - \left(g^{\mu\nu} - \frac{Q^\mu Q^\nu}{Q^2} \right) \left[(Q^2 - \hat{s})^2 + (Q^2 - \hat{u})^2 \right] \right\}. \end{aligned} \quad (11)$$

With our choice of the helicity frame, we have the polarization vector for a longitudinally polarized virtual photon as [9]

$$\epsilon_L^\mu(Q) = -\frac{Q}{\sqrt{(P \cdot Q)^2 - Q^2 S}} \left[P^\mu - \frac{P \cdot Q}{Q^2} Q^\mu \right] \quad (12)$$

where $P^\mu \equiv P_A^\mu + P_B^\mu$. The polarization vector $\epsilon_L^\mu(Q)$ in Eq. (12) reduces to the unit vector \hat{z}^μ in the photon's rest frame [9]. By contracting the hadronic tensors in Eqs. (10) and (11) with the polarization tensor $\epsilon_L^{*\mu}(Q)\epsilon_L^\nu(Q)$ (or $-g^{\mu\nu}$), one can easily obtain the matrix element squares for producing longitudinally polarized (or unpolarized) virtual photons. Because of the ambiguities in choosing the x - or y -axis, the functional forms of the transverse polarization vectors $\epsilon_{T_i}^\mu(Q)$ with $i = 1, 2$ are not unique. One can derive the matrix element squares for producing a transversely polarized virtual photon by using the polarization tensor $\frac{1}{2} \sum_{i=1,2} \epsilon_{T_i}^{*\mu}(Q)\epsilon_{T_i}^\nu(Q)$ [9], or by taking the one half of the difference in the matrix element squares for unpolarized and longitudinally polarized virtual photons.

For deriving our numerical results, we use CTEQ5M parton distributions [15] and the virtual photon fragmentation functions from Ref. [4]. We set the renormalization, factorization, and fragmentation scales equal to $\mu = \kappa \sqrt{Q^2 + Q_T^2}$ with κ a constant of order one.

With the large Q_T setting up the scale of hard collision, the QCD factorization formula in Eq. (5) should be valid for Q^2 as small as a few GeV^2 . Since we are interested in the kinematics similar to J/ψ production at high Q_T , we choose $Q = 2$ and 5 GeV in following plots.

In Fig. 3, we plot the polarization parameter α_{DY} in the proton-antiproton collisions as a function of Q_T at $\sqrt{S} = 2 \text{ TeV}$ (new Tevatron energy). The dashed lines represent the α_{DY} with only the leading order Y -term derived from Eqs. (10) and (11). The solid lines are equal to the α_{DY} with both the Y -term and the resummed contributions. In Fig. 4, we plot the same polarization parameter α_{DY} in the proton-proton collisions at $\sqrt{S} = 500 \text{ GeV}$ (the RHIC energy). The constant $\kappa = 1$ in both Figs. 3 and 4. By varying the κ from $1/2$ to 2 , we find that the cross sections are not very sensitive to the choice of κ [11], and the numerical values of the α_{DY} are relatively stable.

From Figs. 3 and 4, we see that the virtual photons in Drell-Yan massive lepton-pair production are likely to be transversely polarized at high Q_T . The resummed contributions provide significant corrections to the numerical values of the polarization parameter α_{DY} , and they are more important at the low Q^2 . As pointed out in Ref. [4], the longitudinally polarized virtual photon fragmentation functions dominate the threshold region. Consequently, we expect that the resummed contributions to the cross sections of longitudinally polarized virtual photons are more important in low Q_T than high Q_T region, which are clearly evident in the difference between the solid and dashed lines. Since the perturbatively calculable partonic hard parts in the modified QCD factorization formula in Eq. (5) are evaluated at a single hard scale Q_T without large logarithms, high order corrections should not alter the polarization parameter α_{DY} dramatically [11]. We therefore conclude that the virtual photons of invariant mass Q in Drell-Yan massive lepton-pair production are likely to be transversely polarized when $Q_T \gg Q$.

When $Q_T \gg M_{J/\psi}$, hadronic J/ψ and the virtual photon production have a lot in common. The key difference is the difference between their respective fragmentation functions, as shown in Fig. 1. As argued in Ref. [4], the virtual photon fragmentation functions are completely perturbative, while the parton to J/ψ fragmentation functions involve final-state radiations and soft interactions [6]. As shown in Fig. 1, the fragmentation functions from a

parton d to a physical J/ψ can be approximated by the fragmentation functions to a virtual gluon of invariant mass Q , which immediately decay into a $c\bar{c}$ -pair, convoluted with a transition from the $c\bar{c}$ -pair to a physical J/ψ [6]. The first part of the fragmentation functions for a parton to a virtual gluon should be perturbatively calculable. Since a gluon can directly couple to a gluon, the evolution equations for the virtual gluon fragmentation functions can have a nonvanish leading order inhomogeneous evolution kernel for a gluon to a virtual gluon of mass Q . Consequently, the quark-to-virtual-gluon and gluon-to-virtual-gluon fragmentation functions should be equally important, while the gluon-to-virtual-photon fragmentation functions are much smaller than quark-to-virtual-photon fragmentation functions [4]. However, this difference should not have too much effect on the ratio of the transverse and longitudinal polarizations. We then expect that at high Q_T , the virtual gluon, so as the $c\bar{c}$ -pair immediately produced from the decay of the virtual gluon, are more likely to be transversely polarized.

The produced charm and anticharm quark pair of invariant mass Q can in principle radiate gluons and have soft interactions with other partons in the collisions. However, due to the heavy quark mass, such final-state interactions during the transition from the produced $c\bar{c}$ -pair to a physical J/ψ meson are not expected to significantly change the polarization [2]. If the formation from the $c\bar{c}$ -pair of invariant mass Q to a physical J/ψ meson does not change the polarization, one can expect the polarization of the J/ψ mesons produced at high Q_T to be similar to the polarization of the virtual photon in Drell-Yan massive lepton-pair production at the same kinematics. For example, the non-relativistic QCD (NRQCD) model of J/ψ production precisely predicts the J/ψ mesons to be transversely polarized at large Q_T [2], which is not consistent with recent Fermilab data [1].

Since J/ψ production at high Q_T involves both perturbative and nonperturbative physics, we have to find answers to the following two questions in order to resolve this inconsistency. The first question is about the reliability of the perturbative calculations for the relevant kinematics. The second question is about the separation between perturbative and nonperturbative physics. These two questions can be summarized into one: if there is a reliable QCD factorization for hadronic J/ψ production [6]. The measurements of the virtual photon polarization in Drell-Yan massive lepton-pair production can help us to answer the first question. Since the QCD factorization for Drell-Yan massive lepton-pair production is expected to be valid when $Q_T \gg Q$ [11], the measurements of the polarization parameter α_{DY} provide excellent tests of the reliability of perturbative calculations for the kinematic region similar to J/ψ production. Although the measurements of the α_{DY} cannot test the final-state effect of J/ψ production or answer the second question, the measured difference between α_{DY} and $\alpha_{J/\psi}$ can help us to isolate the role of the final-state effect, and to narrow the questions about the J/ψ production.

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FIGURES

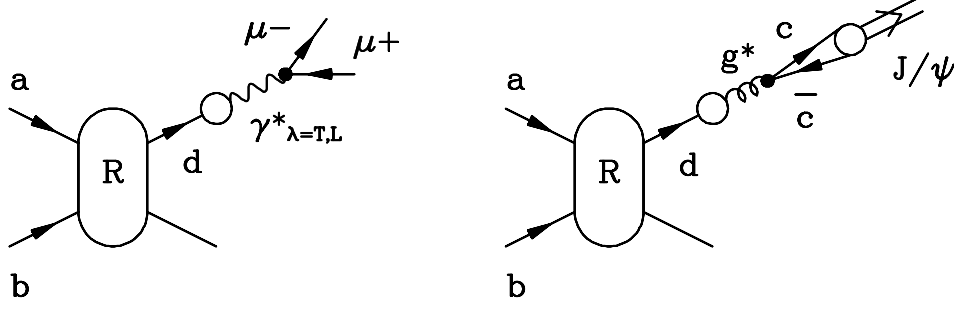


FIG. 1. Sketch for Drell-Yan massive lepton-pair and J/ψ production via parton fragmentation.

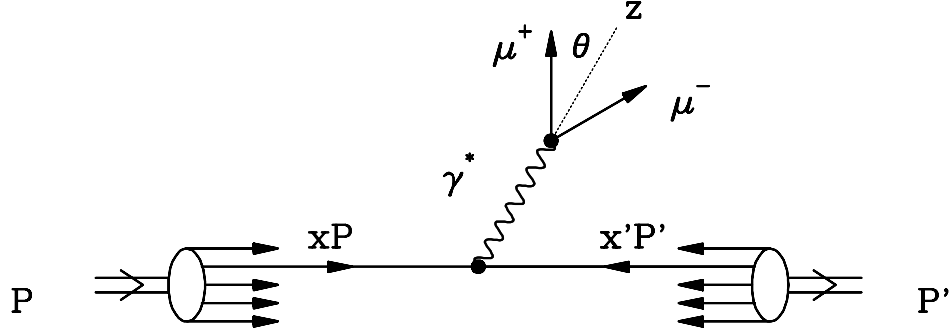


FIG. 2. Definition of the angle θ in Eq. (1) for Drell-Yan massive lepton-pair production.

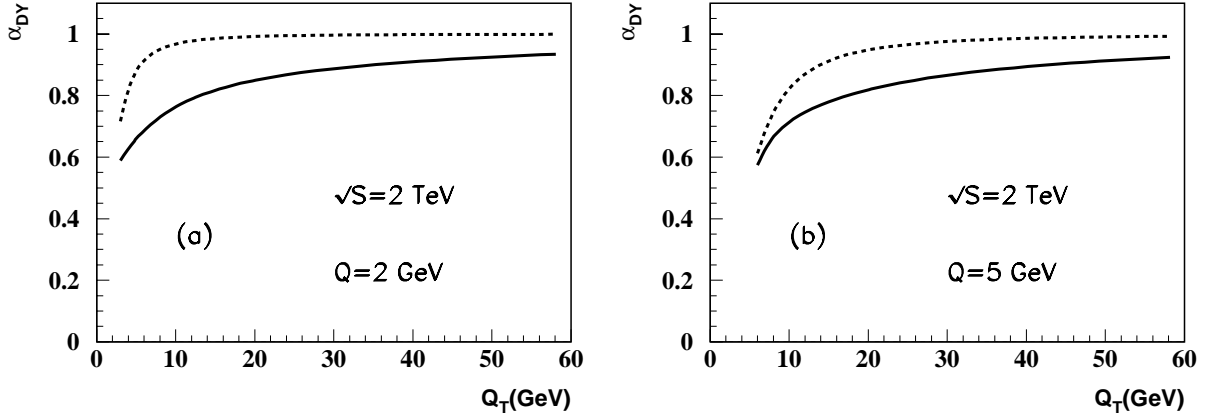


FIG. 3. α_{DY} as a function of Q_T in proton-antiproton collisions at $\sqrt{S} = 2$ TeV and $Q = 2$ GeV (a) and $Q = 5$ GeV (b). Dashed lines correspond to the Y-term only, and solid lines include both the Y-term and the resummed contributions.

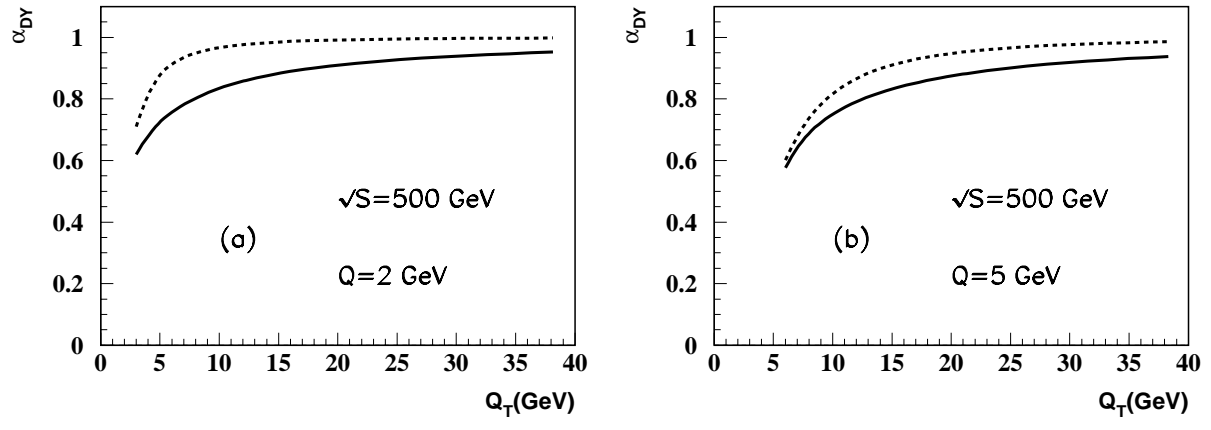


FIG. 4. α_{DY} as a function of Q_T in proton-proton collisions at $\sqrt{S} = 500$ GeV and $Q = 2$ GeV (a) and $Q = 5$ GeV (b). Dashed lines correspond to the Y -term only, and solid lines include both the Y -term and the resummed contributions.

